Dynamical Symmetry Breaking with Vector Bosons

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Abstract

An alternative nonrenormalizable low energy effective model of electroweak symmetry breaking is proposed. In the standard model of electroweak interactions the Higgs doublet is replaced by a complex vector doublet and a real vector singlet. The gauge symmetry is broken dynamically by a mixed condensate of the doublet and singlet vector fields. Gauge fields get their usual standard model masses by condensation. The new vector matter fields become massive by their gauge invariant selfcouplings and expected to have masses of few hundred GeV. Fermions are assigned to the gauge group in the usual manner. Fermion masses are coming from a gauge invariant fermion-vector field interaction by a mixed condensat, the Kobayashi-Maskawa description is unchanged. Perturbative unitarity estimates show that the model is valid up to 2-3 TeV. It is shown that from the new matter fields a large number of spin-one particle pairs is expected at future high energy e^+e^- linear colliders of 500-1500 GeV. The inclusive production cross section of new particle pairs is presented for hadron colliders, while at the Tevatron the new particle production is too low, at the LHC the yield is large.

1 Introduction

The Standard Model of particle physics successfully describes known collider experiments. With the 10 year old discovery of the heavy top quark and the more recent identification of the tau lepton the only missing ingredient of the Standard Model is the elementary Higgs field. In the minimal Standard Model a weak doublet (hypercharge Y=1) scalar field is postulated with an *ad hoc* scalar potential to trigger electroweak symmetry breaking. Three Goldstone Bosons are eaten up by the $W^{\pm}, Z$ and the remaining single CP-even neutral Higgs scalar has evaded the experimental discovery so far. There is only a lower bound from LEP2 experiment $M_H > 114.5$ GeV and an upper bound from electroweak precision tests, which was raised considerably to $M_H <$251 GeV after the refined measurement of the top

quark mass [1]. Beside the missing experimental discovery, theories with elementary scalars are burdened with theoretical problems, like triviality and the most severe gauge hierarchy problem. Elementary scalars are unstable against radiative corrections and without fine tuning the Standard Model must be cut off at 1-2 TeV. We need experiments at the TeV scale to reveal the true nature of electroweak symmetry breaking.

There are basicly two ways to solve these problems in particle physics, either impose new symmetries to protect the scalars or elminate elementary scalars from the theory. The traditional protector is supersymmetry leading to the Minimal Supersymmetric Standard Model, which is very attractive considering the radiatively triggered symmetry breaking, ideal dark matter candidates and successful gauge coupling unification in supersymmetric Grand Unified Theories. However supersymmetric theories involve huge parameter space and doubling of all known particles. None of the predicted new superpartners have been found in any of the experiments and supersymmetry starts to lose it's appeal. Recently "little" Higgs models [2] attracted considerable interest solving the "little hierarchy problem" allowing to raise the cutoff of the theory up to 10 TeV without excessive fine tuning. Little Higgs models realize the old idea that the Higgs is a pseudo Goldstone boson of some spontaneously broken global symmetry [3]. Contrary to supersymmetric models divergent fermion (boson) loops cancel fermion (boson) loops. Little Higgs models still require large fine tuning unless they posess custodial symmetry at the price of highly extended gauge groups.

Scenarios without elementary scalars in four dimemsions are based on dynamical symmetry breaking mechanism. One possibility is a symmetry breaking system interacting strongly with the longitudinal weak vector bosons which has been realised by Dobado *et al.* in the DHT model [4] based on a chiral Lagrangian approach. An alternative description of the strongly interacting symmetry breaking system has been proposed in the BESS model [5] through nonlinear realisations. Top quark condensation has also been suggested for describing the electroweak symmetry breaking [6] leading to several interesting studies [7]. Electroweak symmetry breaking caused by the condensation of a vector field was studied, too [8]. Condensation of vector bosons in different scenarios was considered in the literature. Condensation of the weak gauge bosons $W^{\pm}$ were studied by Linde in extreme dense fermionic matter [10], Ambjorn and Olesen investigated the $W^{\pm}$ condensation in strong external magnetic field [11]. Introduction of new global symmetry and the role of a chemical potential was studied even for the Electroweak Symmetry [12] .

Extra dimensional models were recently proposed to solve the hierarchy problem via changing the high energy (small distance) behaviour of gravity [13, 14]. Concerning symmetry breaking mechanisms new studies went back to the idea of Manton [15] and Hosotani [16] in which the Higgs field is the extra component of a higher dimensional gauge field [17]. Little Higgs models were motivated by these ideas though the simplest models [18] are viable

in four dimensions. The latest idea is breaking the electroweak symmetry without a Higgs [19]. In five dimensions boundary conditions break the original symmetries and the Kaluza Klein modes of the gauge bosons play the role of the standard Higgs scalar, e.g. successfully unitarize the weak gauge boson scattering amplitudes. Higgsless models can be thought as a gravity dual of walking technicolor models with the advantage of weakly coupled regimes allowing perturbative calculations [20]. In the previous examples we have seen that vector bosons can play a role in electroweak symmetry breaking.

In this chapter we present a recently proposed alternative model of electroweak symmetry breaking [9]. Consider the usual Lagrangian of the standard model of electroweak interactions but instead of the scalar doublet two new matter fields are introduced. One of them is a Y = 1, T = 1/2 doublet of complex vector fields

$$B_\mu = \begin{pmatrix} B_\mu^{(+)} \\ B_\mu^{(0)} \end{pmatrix}, \tag{1}$$

the other is a real Y = 0, T = 0 vector field

$$C_\mu. \tag{2}$$

This extends our recent model [8, 21] where only the doublet B_μ was present with the condensation of $B_\mu^{(0)}$. Consequently, we are able to describe a more complete symmetry breaking and to generate fermion masses from a gauge invariant interaction Lagrangian while the mass ratio of the new particles does not become fixed. The key point is the introduction of a mixed $B_\mu - C_\mu$ condensate together with suitable gauge invariant interactions of the new matter fields. This leads to nonvanishing standard model particle masses, as well as B, C particle masses. It turns out that altogether three condensate emerge but only one combination of theirs is fixed by the Fermi coupling constant. The model should be considered as a low energy effective one. Based on recent experience [9, 22], probably it has a few TeV cutoff scale. Its' new particle content is a charged vector boson pair and three neutral vector bosons. As is shown, these can be pair produced in e^+e^- annihilation and hadron colliders, and at future linear colliders of 500-1500 GeV and the LHC they can provide a large number of events.

2 The model

To build the model, in the Lagrangian of the standard model the interactions of the scalar doublet are replaced by the gauge invariant Lagrangian

$$
\begin{aligned}
L_{BC} &= -\frac{1}{2}\overline{\left(D_\mu B_\nu - D_\nu B_\mu\right)}\left(D^\mu B^\nu - D^\nu B^\mu\right) - \\
&\quad -\frac{1}{2}\left(\partial_\mu C_\nu - \partial_\nu C_\mu\right)\left(\partial^\mu C^\nu - \partial^\nu C^\mu\right) - V(B,C),
\end{aligned}
\tag{3}
$$

where D_μ is the covariant derivative, $g_{\mu\nu} = + - --$, and for the potential $V(B,C)$ we assume

$$
V(B,C) = \lambda_1 \left(\overline{B}_\nu B^\nu\right)^2 + \lambda_2 \left(C_\nu C^\nu\right)^2 + \lambda_3 \overline{B}_\nu B^\nu C_\mu C^\mu, \tag{4}
$$

depending only on B-, C- lengths. Other quartic terms would not change the argument. $\lambda_{1,2,3}$ are real and from positivity

$$
\lambda_1 > 0, \quad \lambda_2 > 0, \quad 4\lambda_1\lambda_2 > \lambda_3^2. \tag{5}
$$

Mass terms to be generated are not introduced explicitely in (4). Fermion-BC interactions are introduced later on.

To break the gauge symmetry, we assume a nonvanishing mixed condensate in the vacuum ,

$$
\left\langle \overline{B}_\mu C_\nu \right\rangle = g_{\mu\nu}\ (0,d), \tag{6}
$$

where the left-hand side could be rotated into (0,d), $d \neq 0$, respecting also electric charge conservation and defining the neutral and charged components in (1). By $U_Y(1)$ d can be chosen real. A real d respects also combined TCP and C symmetries. By TCP-invariance (6) equals $\left\langle C_\nu \overline{B}_\mu \right\rangle$. It follows from (6) that the only nonvanishing mixed condensate is

$$
\left\langle B_{1\mu} C_\nu \right\rangle = \sqrt{2} g_{\mu\nu} d \tag{7}
$$

with

$$
B_\mu^{(0)} = \frac{1}{\sqrt{2}}\left(B_{1\mu} + iB_{2\mu}\right), \tag{8}
$$

where $B_{j\mu}$ is real. Once there exists the mixed condensate, B and C are assumed to condense separately

$$
\begin{aligned}
\left\langle \overline{B}_\mu B_\nu \right\rangle &= g_{\mu\nu} k_1, \quad k_1 \neq 0, \\
\left\langle C_\mu C_\nu \right\rangle &= g_{\mu\nu} k_3, \quad k_3 \neq 0.
\end{aligned}
\tag{9}
$$

Once the condensate in (6) defines the broken generators, (9) respects the correct electroweak

symmetry breaking pattern, k_1 must originate from $B_\mu^{(0)}$ condensation,

$$\begin{aligned} \left\langle B_\mu^{(+)\dagger} B_\nu^{(+)} \right\rangle &= 0, \\ \left\langle B_\mu^{(0)\dagger} B_\nu^{(0)} \right\rangle &= g_{\mu\nu} k_1. \end{aligned} \tag{10}$$

(10) reproduces the pattern of gauge particle masses [8]. All the condensates linear in $B_\mu^{(+)}$ vanish by charge conservation. Finally, we assume in advance, that

$$\left\langle B_\mu^{(0)} B_\nu^{(0)} \right\rangle = \left\langle B_\mu^{(0)\dagger} B_\nu^{(0)\dagger} \right\rangle = g_{\mu\nu}\ k_2. \tag{11}$$

The point is that in general $B_{1\mu}$ and $B_{2\mu}$ belong to different masses, so that $k_2 \neq 0$. $k_{1,2,3}$ are real and $k_1 < 0$, $k_3 < 0$, as shown by physical particle masses and simple models. The condensates are of nonperturbative origin caused by the strong interaction V(B,C). Among them only k_1 is fixed by contemporary phenomenology.

3 Boson and fermion masses

Mass terms are obtained in the linearized form of L_{BC} via condensates. The $W^\pm$ mass is determined by the total B-condensate, while the two neutral gauge field combinations are proportional to $B_\mu^{(+)\dagger} B_\nu^{(+)}$ and $B_\mu^{(0)\dagger} B_\nu^{(0)}$, respectively. Therefore, (10) yields

$$m_{\mathrm{photon}} = 0,\ m_W = \frac{g}{2}\sqrt{-6k_1},\ m_Z = \frac{g}{2\cos\theta_W}\sqrt{-6k_1}. \tag{12}$$

Low energy phenomenology gives

$$k_1 = -\left(6\sqrt{2}G_F\right)^{-1},\ (-6k_1)^{1/2} = 246 \text{ GeV}. \tag{13}$$

$B^\pm$ and B_2 get the following masses

$$\begin{aligned} m_\pm^2 &= -8\lambda_1 k_1 - 4\lambda_3 k_3, \\ m_{B_2}^2 &= -10\lambda_1 k_1 + 2\lambda_1 k_2 - 4\lambda_3 k_3 = m_\pm^2 + 2\lambda_1(k_2 - k_1). \end{aligned} \tag{14}$$

For $\lambda_3, -k_1, -k_3 > 0$, $m_{B_2}^2 > m_\pm^2 > 0$ since $k_2 > k_1$. The $B_1 - C$ sector is slightly more complicated, here one arrives at the following bilinear combinations in the potential for $B_{1\mu}, C_\mu$

$$V(B,C) \to -\frac{m_1^2}{2} B_{1\mu} B^{1\mu} - \frac{m_2^2}{2} C_\nu C^\nu - m_3^2 B_{1\mu} C^\mu, \tag{15}$$

with

$$
\begin{aligned}
-m_1^2 &= 10\lambda_1 k_1 + 2\lambda_1 k_2 + 4\lambda_3 k_3 = -m_{B_2}^2 + 4\lambda_1 k_2, \\
-m_2^2 &= 24\lambda_2 k_3 + 8\lambda_3 k_1, \\
-m_3^2 &= 4\sqrt{2}\lambda_3 d.
\end{aligned} \tag{16}
$$

Here $m_1^2 > 0$ being $k_1 + k_2 < 0$; $m_2^2 > 0$ and $m_3^2 \lesseqgtr 0$. A positive potential in (14) requires

$$
m_1^2, m_2^2 > 0, \qquad m_1^2 m_2^2 > m_3^4. \tag{17}
$$

(15) shows that $B_{1\mu}$ and $C_\mu$are nonphysical fields, the mass eigenstates are defined by

$$
\begin{aligned}
B_{f\mu} &= cB_{1\mu} + sC_\mu, \\
C_{f\mu} &= -sB_{1\mu} + cC_\mu,
\end{aligned} \tag{18}
$$

where $c = \cos\phi,\ s = \sin\phi$, ϕ denotes the mixing angle defined by

$$
\frac{1}{2}\sin 2\phi(m_1^2 - m_2^2) = \cos 2\phi m_3^2. \tag{19}
$$

The physical masses are

$$
\begin{aligned}
m_{B_f}^2 &= c^2 m_1^2 + s^2 m_2^2 + 2csm_3^2, \\
m_{C_f}^2 &= s^2 m_1^1 + c^2 m_2^2 - 2csm_3^2,
\end{aligned} \tag{20}
$$

whence

$$
2m_{B_f}^2 = m_1^2 + m_2^2 + \frac{m_1^2 - m_2^2}{\cos 2\phi}. \tag{21}
$$

$$
2m_{C_f}^2 = m_1^2 + m_2^2 - \frac{m_1^2 - m_2^2}{\cos 2\phi} \tag{22}
$$

For $(m_1^2 - m_2^2)/\cos 2\phi > 0$ (< 0) $m_{B_f}^2 > m_{C_f}^2 > 0$ $(m_{C_f}^2 > m_{B_f}^2 > 0)$. At vanishing mixing, $m_3^2 = 0$, $B_{1\mu}$ and C_μ become independent having the masses m_1 and m_2; taking $k_2 = 0$ and omitting C_μ we recover the model of Ref.5. k_2 shifts the real component field masses from the mass of the imaginary part $B_{2\mu}$.

The particle spectrum of the B-C sector consists of the spin-one $B^\pm$ and the three neutral spin-one particles B_2, B_f, C_f. Their masses are rather weakly restricted. Beside the gauge coupling constants and $\lambda_1, \lambda_2, \lambda_3$, the model has three basic condensates $\langle V_{i\mu} V_{i\nu}\rangle$,

$V_{i\mu} = B_{2\mu}, B_{f\mu}, C_{f\mu}$. k_1, k_2, k_3, d condensates are built up from these as follows

$$\begin{aligned}
g_{\mu\nu} d &= \frac{1}{\sqrt{2}} cs \left(\langle B_{f\mu} B_{f\nu} \rangle - \langle C_{f\mu} C_{f\nu} \rangle \right), \\
g_{\mu\nu} k_1 &= \frac{1}{2} \left\{ c^2 \langle B_{f\mu} B_{f\nu} \rangle + s^2 \langle C_{f\mu} C_{f\nu} \rangle + \langle B_{2\mu} B_{2\nu} \rangle \right\}, \\
g_{\mu\nu} k_2 &= \frac{1}{2} \left\{ c^2 \langle B_{f\mu} B_{f\nu} \rangle + s^2 \langle C_{f\mu} C_{f\nu} \rangle - \langle B_{2\mu} B_{2\nu} \rangle \right\}, \\
g_{\mu\nu} k_3 &= s^2 \langle B_{f\mu} B_{f\nu} \rangle + c^2 \langle C_{f\mu} C_{f\nu} \rangle .
\end{aligned} \tag{23}$$

From (23) d can be written as

$$2\sqrt{2} \cot 2\phi \, d = k_1 + k_2 - k_3. \tag{24}$$

Turning to the dynamical fermion mass generation, we add to the gauge vector and matter vector field Lagrangians, the usual fermion-gauge vector Lagrangian, as well as a new gauge invariant piece responsible for the fermion-matter vector field interactions and in usual notation this is (for quarks)

$$g^u_{ij} \overline{\psi}_{iL} u_{jR} B^C_\nu C^\nu + g^d_{ij} \overline{\psi}_{iL} d_{jR} B_\nu c^\nu + h.c., \tag{25}$$

$$\psi_{iL} = \begin{pmatrix} u_i \\ d_i \end{pmatrix}_L, \quad B^C_\nu = \begin{pmatrix} B^{(0)\dagger}_\nu \\ -B^{(+)\dagger}_\nu \end{pmatrix}.$$

Clearly the mixed condensate provides fermion masses and also the Kobayashi-Maskawa description is unchanged. A typical fermion mass is

$$m_f = -4 g_f d \tag{26}$$

and only $g_f d$ becomes fixed but $m_{f1}/m_{f2} = g_{f1}/g_{f2}$ as usual. If d is about $k_1 \simeq G_F^{-1}$, then g_f is a factor of $G_F^{1/2}$ weaker than the approximate standard model value $G_F^{1/2}$.

4 Interactions of the new vector bosons

The new physical fields $B_{f\mu}, C_{f\mu}, B_{2\mu}, B^+_\mu$ have various interactions. Four-boson self-interactions can be read off from the V(B,C) potential (4), the coupling are all proportional to some unknown λ_i.

The new vector bosons interact with the standard fermions via the Yukawa interactions (25), but the coupling strength is expected to be weaker than the Standard Model Higgs-fermion couplings.

The most important interactions of the new particles from the point of view of phenomenology are the one with the gauge bosons. The source of these interaction is the

B_μ covariant derivative terms in (3). The three particle interaction all have a derivative coupling, before the B-C mass diagonalization they are

$$\begin{aligned} L_{int}^{(3)} = \quad & -ie\left(\partial_\mu B_\nu^{(-)} - \partial_\nu B_\mu^{(-)} B^{(+)\nu}\right) A^\mu + \\ & +ie\cot 2\theta_w \left(\partial_\mu B_\nu^{(-)} - \partial_\nu B_\mu^{(-)}\right) B^{(+)\nu} Z^\mu - \\ & -i\tfrac{\sqrt{2}}{2}\tfrac{e}{\sin\theta_w}\left(\partial_\mu B_\nu^{(-)} - \partial_\nu B_\mu^{(-)}\right) B^{(0)+\nu} W^{+\mu} - \\ & -i\tfrac{\sqrt{2}}{2}\tfrac{e}{\sin\theta_w}\left(\partial_\mu B_\nu^{(0)+} - \partial_\nu B_\mu^{(0)+}\right) B^{(+)\nu} W^{-\mu} - \\ & -i\tfrac{e}{\sin 2\theta_w}\left(\partial_\mu B_\nu^{(0)+} - \partial_\nu B_\mu^{(0)+}\right) B^{(0)\nu} Z^\mu \qquad +h.c. \end{aligned} \tag{27}$$

In terms of the physical fields the neutral boson interactions change considerably, as an example we give the $Z - B_f - B_2$ coupling,

$$L_{int,2} = \frac{g}{2\cos\theta_W}\cos\phi\cdot Z_\mu\left[B_{f\nu}(\partial^\mu B_2^\nu - \partial^\nu B_2^\mu) - B_{2\nu}(\partial^\mu B_f^\nu - \partial^\nu B_f^\mu)\right]. \tag{28}$$

There exist several B_iB_jVV- type four particle couplings with gauge bosons with $V = \gamma, W^\pm, Z$ and $B_{i\mu} = B_{f\mu}, C_{f\mu}, B_{2\mu}, B_\mu^+$, the coupling strength is $\sim g^2$ multiplied with mixing angles (e.g. $\cos\theta_W, \sin\phi$). We remind here that there is no mixing between the gauge and new vector bosons. As an example we give the VVB^+B^-couplings

$$\begin{aligned} L_{int,+}^{(4)} \quad = \quad & \left(g^{\mu\alpha}g^{\nu\beta} - g^{\mu\nu}g^{\alpha\beta}\right)\left[e^2 A_\nu A_\mu B_\alpha^{(-)} B_\beta^{(+)} + \frac{g^2}{2} W_\mu^{(-)} W_\mu^{(+)} B_\alpha^{(-)} B_\beta^{(+)} + \right. \\ & \left. + e^2\cot^2 2\theta_W Z_\mu Z_\nu B_\alpha^{(-)} B_\beta^{(+)} - e^2\cot 2\theta_W\left(A_\mu Z_\nu B_\alpha^{(-)} B_\beta^{(+)} + Z_\mu A_\nu B_\alpha^{(-)} B_\beta^{(+)}\right)\right] \end{aligned} \tag{29}$$

All the interactions have even number of new B, C particles, the single production of the new particles are not allowed in this model. The new particles can only decay to each other determined by their mass relation making them difficult to discover in particle physics experiment and the lightest one is stable. At the same time if the lightest new vector boson is neutral then it will be an ideal candidate for a selfinteracting dark matter. Generally all the couplings are weaker than the relevant couplings of one standard Higgs boson and the most important ones are the three boson couplings in (27,28). In what follows we investigate the implications of the interactions.

5 Unitarity constraints

In this section we apply tree-level partial wave unitarity to two-body scatterings of longitudinal gauge and B,C bosons following the reasoning of [23], where perturbative unitarity has been employed to constrain the Standard Model Higgs mass. With the only experimental

input of the muon decay constant G_F the perturbative upper bound of appoximately 1 TeV emerged for the Standard Model Higgs mass. Perturbative unitarity is a powerful tool, it can be used to build up the bosonic sector of the Standard Model and it was essential to build higssless models of electroweak symmetry breaking in extra dimensional field theories. Perturbative unitarity shows that the scale of the model is approximately 2.5 TeV and the new particle masses are bounded from below, the lower bound increasing with a growing Λ [22].

In the vector condensate model there exist many elastic $VV \to B_i B_j$ and $B_i V \to B_i V$ type processes, with $V = \gamma, W^{\pm}, Z$ and $B_i = B_f, C_f, B_2, B^+$. We consider these processes for longitudinally polarized external particles, as these are the most dangerous one to ruin the high energy behaviour, in the high energy limit

$$\epsilon_\mu(k) \sim \frac{k_\mu}{M}, \tag{30}$$

where M is the mass of the outgoing/incoming vector particle. We calculate then the $J = 0$ partial-wave amplitudes, a_0, from contact and one-particle exchange graphs. Unitarity requires $|Re\, a_0| \leq \frac{1}{2}$.

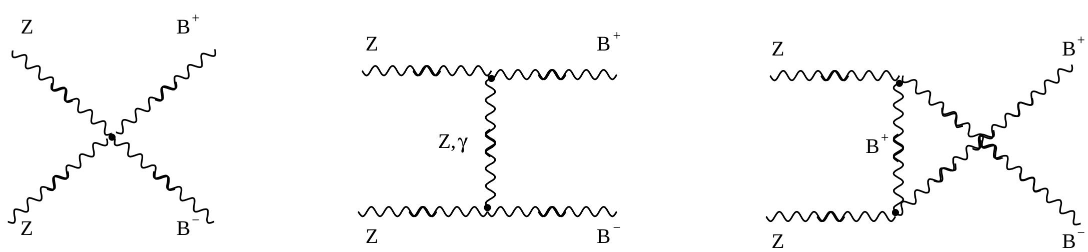

Figure 1: Tree level Feynman graphs for $B^+B^- \to B^+B^-$.

The scattering processes provide vastly different constraints on the parameters of the theory. First choose a process with known coupling constants. In the process $ZZ \to B^+B^-$ all the couplings are proportional to the weak coupling constant ($e \cot 2\theta_W$). There are three graphs shown in Fig 1. contributing to the elastic scattering, the contact graph with the four particle vertex and t- and u-channel graphs. In the $s \gg m_+, m_Z$ limit applying

(30) the contributions are the following ($m_+ = m_{B^+}$)

$$\begin{aligned} T_c &= e^2 \frac{\cot^2 2\theta_W}{m_Z^2 m_+^2} \left(\frac{s^2}{2} - \frac{t^2}{4} - \frac{u^2}{4} \right), \\ T_t &= e^2 \frac{\cot^2 2\theta_W}{m_Z^2 m_+^2} \left(\frac{1}{4} t(s-u) \right), \\ T_u &= e^2 \frac{\cot^2 2\theta_W}{m_Z^2 m_+^2} \left(\frac{1}{4} u(s-t) \right). \end{aligned}$$

The sum of the three amplitude growing with the energy vanishes

$$T_c + T_t + T_u = 0,$$

reflecting that the B^+B^-Z interactions can in principle originate from a renormalizable interactions and the high energy behaviour of the process is modest. A detailed calculation with the general polarizations give for the s-wave amplitude

$$|a_0| = \frac{e^2 \cot^2 2\theta_W}{32\pi} \frac{m_Z^2}{m_+^2} + \mathcal{O}(1/s),$$

giving $m_+ \geq 3$GeV and similar weak bounds emerge from other elastic $BV \to BV$ and $BB \to VV$ scatterings.

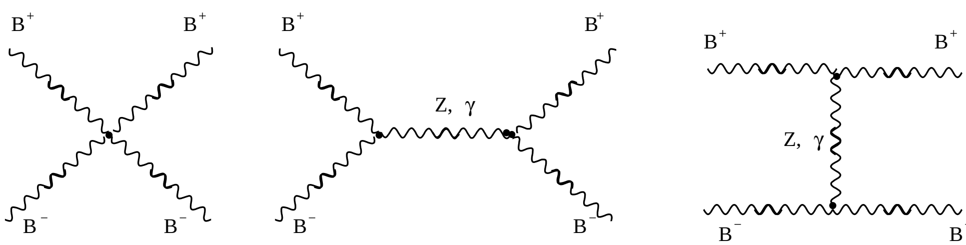

Figure 2: Tree level Feynman graphs for $B^+B^- \to B^+B^-$.

Strong bounds emerge from the really non-renomalizable sector of the theory, the quartic B interactions. Taking $\lambda_3 k_3$ negligible in (14) λ_1 is proportional to $m_+^2 G_F$. Consider the tree level process $B^+B^- \to B^+B^-$ via the contact graph and the Z, γ exchange graphs in the s and t channels (Fig 2.). In the limit $s \gg m_+^2, m_Z^2$ only the contribution of the contact graph is important

$$T_c = \lambda_1 \frac{1}{2m_+^4} \left(t^2 + u^2 \right),$$

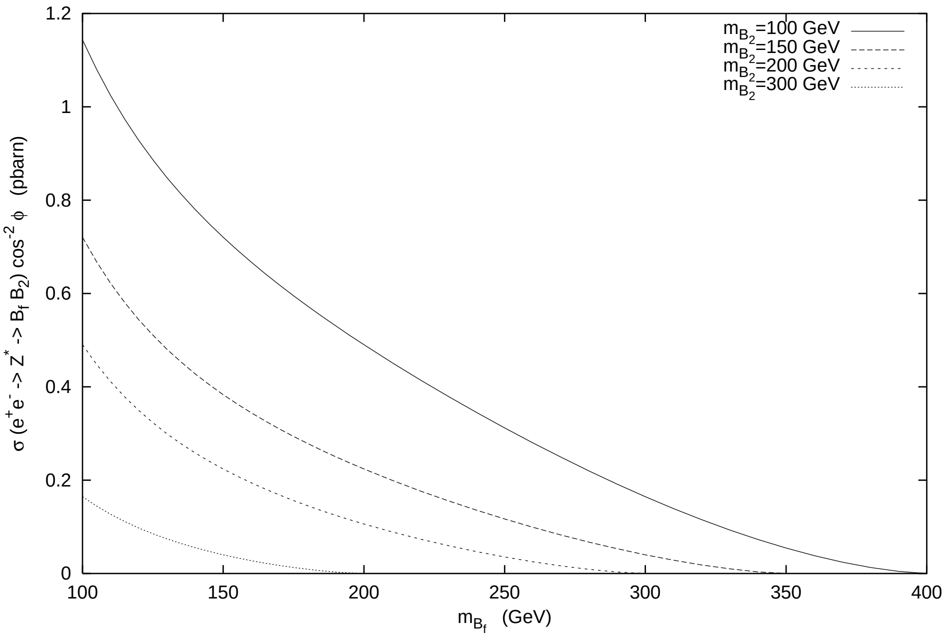

Figure 3: $\cos^{-2}\phi\,\sigma\,(e^+e^- \to B_f B_2)$ vs. m_{B_f}at $\sqrt{s} = 500$ GeV and various m_{B_2}.

yielding for the s-wave amplitude

$$|a_0| = \frac{\sqrt{2}}{64\pi} G_F \frac{s^2}{m_+^2}.$$

The unitarity constraint $|Re\, a_0| \leq \frac{1}{2}$ along with the assumptions that the masses in the model must be smaller than the cutoff provide the scale of the model $\Lambda \leq 2.5$TeV. The bound is similar also for the $B^\pm B^\pm \to B_2 B_2$ scattering. In case of vanishing mixing between B_1 and C a rough interpretation of k_1 with a cutoff free propagator shows similar bounds $\Lambda \leq 2 - 2.6$ TeV depending on the interpretation of k_1 [24]. We conclude that perturbative unitarity estimates suggest that the scale of the model is in the range 2-3 TeV, where a (more) fundamental description takes the role of the Vector Condesate Model.

6 Production in e^+e^-colliders

Direct production of $B_f B_2$ pairs can be studied in high energy e^+e^- colliders, $e^+e^- \to Z^* \to B_f B_2$. Assume in (26) g_{e^-} is very small, then the direct $e^+e^- \to B_f B_2$ can be neglected. From (28) we have for the total cross section

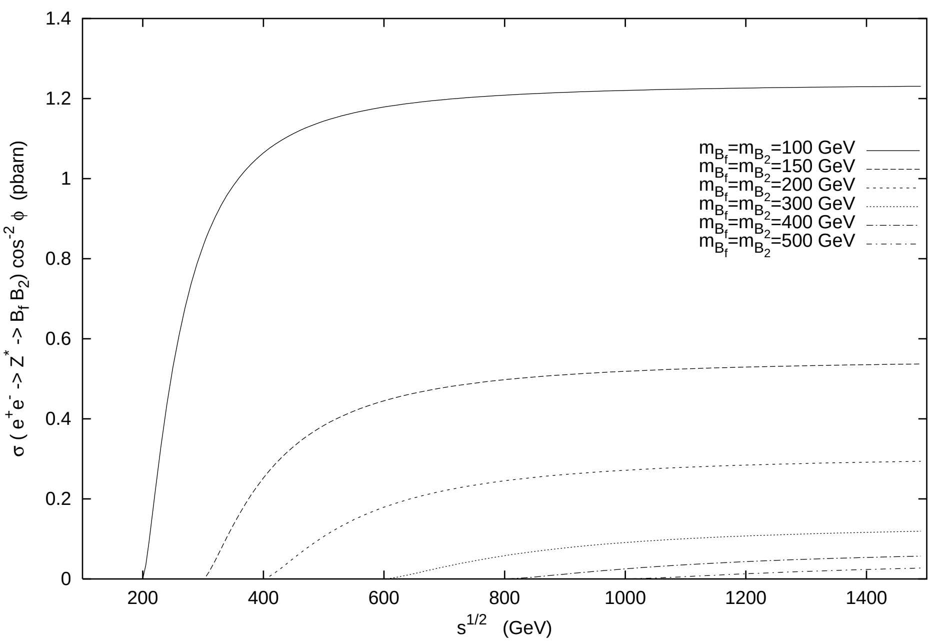

Figure 4: $\cos^{-2}\phi\,\sigma\,(e^+e^- \to B_f B_2)$ vs. $\sqrt{s}$ at various $m_{B_f} = m_{B_2}$.

$$\sigma(e^+e^- \to Z^* \to B_f B_2) = \frac{g^4 \cos^2\phi}{3\cdot 4096 \cos^4\theta_W} \frac{1+(4\sin^2\theta_W - 1)^2}{m_{B_2}^2 m_{B_f}^2 s^2 (s-M_Z^2)^2} \left(s-(m_{B_2}+m_{B_f})^2\right)^{3/2} \cdot$$

$$\cdot \left(s-(m_{B_f}-m_{B_2})^2\right)^{3/2} \left(2s(m_{B_2}^2+m_{B_f}^2)+m_{B_f}^4+m_{B_2}^4+10m_{B_2}^2 m_{B_f}^2\right). \qquad (31)$$

At asymptotic energies σ is proportional to $1/m_{B_2}^2 + 1/m_{B_f}^2$. The mass and energy dependences of σ are shown in Figs. 3,4. For example at $\sqrt{s}=$ 500 GeV and with an integrated luminosity of 10 fb^{-1} 5700, 1900, 530 $B_f B_2$ pairs are expected for $m_{B_f} = m_{B_2}=$ 100, 150, 200 GeV and $\cos^2\phi = 1/2$. At $\sqrt{s} = 1.5$ TeV a higher mass range can be tested, for $\cos^2\phi = 1/2$, a luminosity of 100 fb^{-1} we get the large event numbers 62200, 14500, 5900, 1900, 530 for $m_{B_f} = m_{B_2}=$ 100, 200, 300, 400, 500 GeV. One can show that the B^+B^- production is a factor of $\cos^2 2\theta_W$ smaller than (31) at equal masses and $\cos^2\phi = 1$.

Though we expect a large number of events at the next generation of electron positron colliders the identification of the new particles will be a difficult task. The lightest of B, C is stable and if it is a neutral particle it is ideal dark matter candidate then the same problem expected concerning the detection as in the case of other dark matter particles. The lightest neutral particle only turns up in the missing energy channels. When a heavier new particle

is produced it can decay subsequently into a lighter new particle emitting a single gauge boson or a pair of leptons, the appearance of the standard decay products at a misplaced vertex together with missing energy will trigger the hopeful discovery of the vector bosons. The charged B^+-s can be identified fairly easier via charged tracks and missing energy in the calorimeters.

7 New particles at hadron colliders

In this section we study the production of B–particles at the most energetic hadron colliders, ideal for discoveries, Tevatron and LHC. We show that producing heavy B(C)-particles at LHC is very favourable having a large cross section while at the Tevatron energy the production cross section cannot exceed (0.01-0.02) fb which is far below the discovery limit.

Since fermions are coupled very weakly to B(C)-pairs in the vector condensate model, producing B,C-pairs is expected to be more considerable from virtual γ and Z exchanges, that is we consider the Drell-Yan mechanism [26], $p\overset{(-)}{p} \to B\overline{B} + X$ via quark-antiquark annihilation.

The Drell-Yan cross section for the above hadronic collisions can be written as [26, 27]

$$\sigma(p\overset{(-)}{p} \to B\overline{B} + X) = \int_{\tau_0}^{1} d\tau \int_{\tau}^{1} \frac{dx}{2x} \sum_i \sigma(q_i\overline{q}_i \to B\overline{B}) \cdot \left(f_i^1(x,\hat{s}) f_{\bar{i}}^2(\tau/x,\hat{s}) + f_{\bar{i}}^1(x,\hat{s}) f_i^2(\tau/x,\hat{s})\right), \tag{32}$$

where x and τ/x are the parton momentum fractions, $\hat{s} = \tau s$ is the square of the centre of mass energy of $q_i\bar{q}_i$, s is the same for the hadronic initial state, $f_i^1(x,\hat{s})$ means the number distribution of i quarks in hadron 1 at the scale $\hat{s}$ and the sum runs over the quark flavours u,d,s,c. In the computation the MRS (G) fit program [28] was used for the parton distributions.

The angle integrated, colour averaged annihilation cross section $\sigma(q_i\overline{q}_i \to B^+B^-)$ is calculated to lowest order in the gauge couplings, and QCD corrections are neglected. We hope this approximation shows the order of magnitude of the cross section. We give the result of the charged final state as there is no unknown mixing angle in the result and the identification of $B^\pm$ seems less dificult than for neutral pairs. The B^+B^- pairs appear via $\gamma + Z$ exchange, the relevant interactions are in (27). At the $q_j\overline{q}_j Z$-vertex the usual coupling $ig\gamma_\mu(g_{Vj} + g_{Aj}\gamma_5)$ acts, here

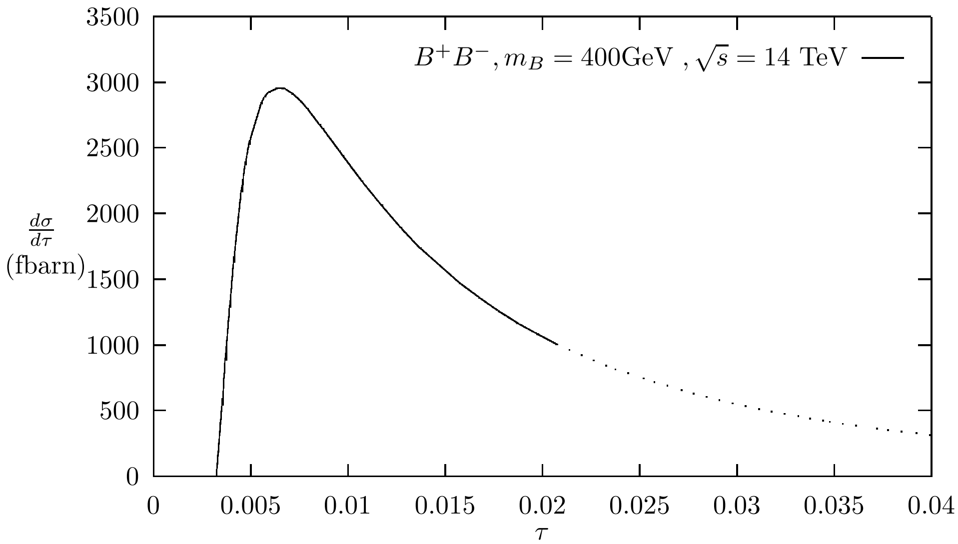

Figure 5: The differential τ-distribution of B^+B^- pairs at LHC, $m_+ = 400$ GeV.

$$\left.\begin{array}{ll} g_{Vj} = & \frac{1}{2} - \frac{4}{3}\sin^2\theta_W, \\ g_{Aj} = & \frac{1}{2} \end{array}\right\} j = u, c$$

$$\left.\begin{array}{ll} g_{Vj} = & -\frac{1}{2} + \frac{2}{3}\sin^2\theta_W, \\ g_{Aj} = & -\frac{1}{2} \end{array}\right\} j = d, s \ . \tag{33}$$

From (27) we get for B^+B^- final states [29]

$$\begin{aligned} \sigma\left(q_i\overline{q}_i \to B^+B^-\right) &= \left[(g_{Vi}^2 + g_{Ai}^2)cos^2 2\theta_W + 2Q_{q_i} g_{Vi} sin^2 2\theta_W cos2\theta_W + 4Q_{q_i}^2 sin^4 2\theta_W\right] \\ & \frac{1}{3}\frac{1}{256\pi}\left(\frac{g}{cos\theta_W}\right)^4\left(1 - \frac{4m_+^2}{\hat{s}}\right)^{3/2}\frac{\hat{s} + 8m_+^2}{4m_+^4} \end{aligned} \tag{34}$$

This is decreasing at high, increasing m_+ and for $\hat{s} \gg 4m_+^2$ it is proportional to $\hat{s}/m_+^4$ reflecting that the Lagrangian (28) is coming from the nonrenormalizable, effective model. The individual terms are due to Z exchange, $\gamma - Z$ interference and γ exchange. The production of neutral particle pairs B_fB_2 (C_fB_2) is only realized via the Z exchange channel includes an undetermined mixing angle factor, $\cos^2\phi$ ($\sin^2\phi$) and is slighty more involved because of two mass parameters.

We have calculated various distributions of B^+B^- pairs for $p\overline{p}$ collisions at $\sqrt{s} = 1.8$

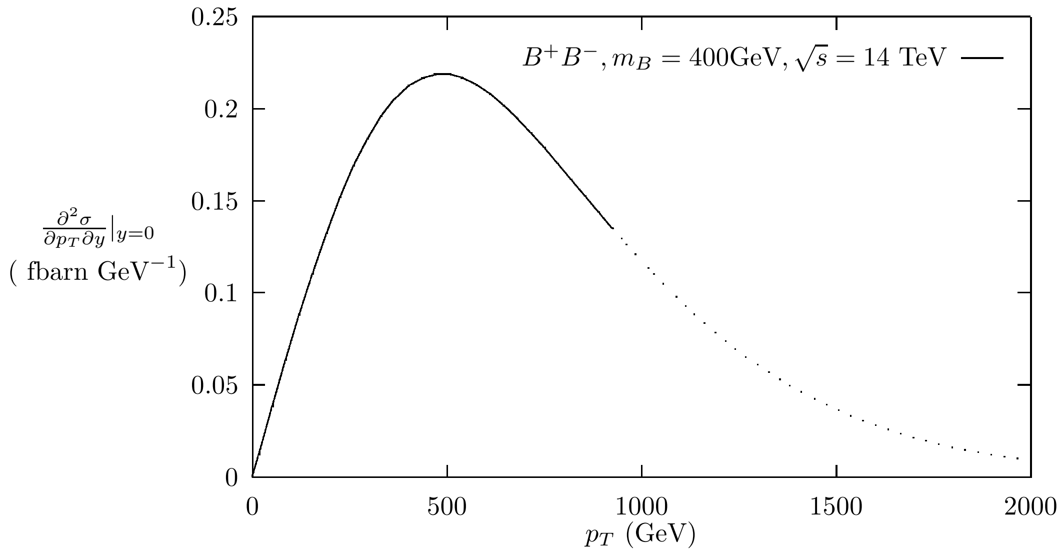

Figure 6: $\frac{\partial^2\sigma}{\partial p_T \partial y}|_{y=0}$ at LHC, $m_+ = 400$ GeV. p_T denotes the transverse momentum of B^+ and y is the rapidity of B^+B^-.

TeV and for pp collisions at $\sqrt{s} = 14$ TeV, assuming $m_+ = 400, 500, 600$ GeV. Typically, at the Tevatron no notable result can be presented, however, increasing the energy to LHC, we get sizable cross sections. Also the yield of B^+B^- is larger than that of neutral pairs $B_f B_2, C_f B_2$. As an example we show in Fig.5. the τ-distribution of B^+B^- pairs at LHC, $m_+ = 400$ GeV. $\frac{d\sigma}{d\tau}$ is sharply peaked after threshold $(4m_+^2/s)$ and decreases for higher invariant masses of B^+B^-. Calculating the total cross section the dotted part of $d\sigma/d\tau$ was not integrated corresponding to a cutoff $\sqrt{\hat{s}} = 2$ TeV at the parton level ($\tau \leq 0.02$). Fig. 6. shows $\frac{\partial^2\sigma}{\partial p_T \partial y}|_{y=0}$ as the function of the transverse momentum p_T of B^+ at vanishing rapidity y of B^+B^- for LHC, $m_+ = 400$ GeV. The start of the dotted curve correspond to $\sqrt{\hat{s}} = 2$ TeV.

For the total cross section (32) we obtain

$$\begin{aligned} \sigma_{\text{Tev}} &= 0.020 \text{ fb} \quad \text{for} \quad B^+B^-, m_+ = 400\text{GeV}, \\ \sigma_{LHC} &= 33.0; 8.5; 2.4\,\text{fb} \quad \text{for} \quad B^+B^-, m_+ = 0.4; 0.5; 0.6\text{TeV}. \end{aligned} \tag{35}$$

From Fig. 5. we can immediately read off the cutoff dependence. For instance, for $m_+ = 400$ GeV, $\sigma_{LHC}(B^+B^-) = 20.5(36.8)$ fb at a cutoff 1.5(2.5) TeV. At an expected integrated luminosity of 10^5pb^{-1} one gets about 3300 B^+B^- pairs of $m_+ = 0.4$ TeV at LHC per annum at a cutoff 2 TeV. The detection of the new particles similar to the case of

electron-positron colliders, though it is even more difficult, charged tracks can be searched for, but the high luminosity provides large number of events.

In this section, we have shown that heavy B–particle pairs have a large inclusive cross section due to $q\overline{q}$ annihilation at the LHC in the few hundred GeV mass range making the detection of B^+ or neutral partners at LHC possible.

8 Conclusion

In this chapter a low energy dynamical symmetry breaking model of electroweak interactions based on matter vector field condensation is introduced. Mass generation is arranged starting from gauge invariant Lagrangians, fermion masses are also coming from the vector condensates. New particles are all spin-one states, one charged pair and three neutral particles having various interactions. The masses of the new particles are assumed to generated also by condensation resulting a nontrivial mixing among the neutral components. The production of the new particles in electron-positron collider was studied and a large number of event is expected at future colliders with the center of mass energy 500-1000 GeV. Tevatron cannot produce enough new particles for observation, but the yield at the LHC expected to be well above the discovery limit for new particles with masses of a few hundred GeV. At present there is only a $\sim$45 GeV direct lower bound from the invisible Z width [24]. Further constraints can arise from the electroweak precision tests of the Standard Model. In the model [8] proposed earlier the precison S,T,U parameters [30] were calculated [8], and found to be in agreement with the latest experimental data. The parameter space of the model presented in this chapter is larger than that of the one in [8], therefore, we expect that the positive result of the S,T parameter analysis can be maintained. It would be interesting and exciting to investigate in details the one loop radiative corrections in our model and to confront with available experiments.

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